hep-ph/9906493
June 1999

The Contributions of Neutral Higgs Bosons to Charmless Nonleptonic B Decays in MSSM

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abstract

We investigate the contributions of neutral Higgs bosons to nonleptonic transition $b \to q\bar{s}s, q = d,\ s$ under the supersymmetric context. Their effects to decay width and CP violation in corresponding exclusive decays are explored. The anomalous dimension matrices of the operators which have to be incorporated to include the contributions of neutral Higgs bosons are given. We find that when $\tan\beta$ is large (say, 50) and neutral Higgs bosons are not too heavy (say, 100 GeV), contributions of neutral Higgs penguin can dominate electroweak penguin contributions, and for some processes, they can greatly modify both decay width and CP asymmetry.

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The systematic study of the nonleptonic decays of B mesons has been led long before the construction of B factories. The CLEO report [1] on the measurement of the charmless two-body hadronic decays revived the theoretical interest in this subject. Among the reported exclusive decay modes, $B \to K\eta'$ received much more attention (for a review, see [2]), and according to [3, 4], four-quark-operator contribution is capable of explaining it under the theoretical framework of SM. For the inclusive decay $B \to X_s\eta'$, it seems that new mechanism or new physics should be introduced in order to explain the unexpected large branching ratio [5].

The effective Hamiltonian method is the fundamental tool to study the charmless non-leptonic decays of B mesons. Starting from a full theory, after integrating out the heavy freedom degree at a certain energy scale, the effective Hamiltonian is obtained. For the standard model (SM), the heavy freedom degrees integrated out at $\mu_0 = m_w$ are top quark, Z boson and W boson. For the SUSY models, in order to be simple, the extra heavy freedom degrees integrated out include all superpartners, though there are superpartners, such as the lightest neutralino, which may be lighter than W boson, and there are superpartners, such as the super-quarks of the first two generations, which may be much heavier than W boson. At this integrating out scale $\mu_0 = m_w$, super-partners can not only contribute to the effective Wilson coefficients, but also introduce new operators which will be shown below in the case of including contributions of neutral Higgs bosons (NHB).

The works on the next-to-leading-order QCD-improved effective Hamiltonian derived from SM were reviewed in [6], where the problem on the scale independence were studied. To present days, our knowledge of the dynamics of the hadronic transitions are limited, and the approach employed in evaluating hadronic matrix elements of operators is factorization. There are two last papers which represent the state of art of the systematic study of charmless hadronic two-body B decays. Ref [3] concentrates on the experimental test on the naive factorization ansatz, having classified 76 decay modes into five classes and analyzed the dependence of these decay modes on the weak mixing matrix elements, form factors, decay constants, QCD scale, and the effective number of color. Ref [4] shows the experimental data of $\rho^0\pi^\pm$ and $\phi K^\pm$ cannot be accommodated simultaneously by treating $N_c^{eff}(LL) = N_c^{eff}(LR)$ and, by assuming that $N_c^{eff}(LL) \simeq 2$ and $N_c^{eff}(LR) \sim 5$, finds the existed experimental data can be explained well.

Nonleptonic decays have also been investigated in SUSY models [7, 8]. Several papers deal with the possible effects of SUSY to hadronic B decays in the SUSY models with R parity violation [9]. The other topic related with the nonleptonic B decays of SUSY models is CP violation, there are a vast amount of literatures [10, 11].

It is well known that for charmless B hadronic decays with $\Delta B = 1$ in which tree level contributions are suppressed by CKM and color symmetry, penguin contributions, among which gluonic ones play major parts, always dominate. In this letter we investigate the total contributions of penguins in MSSM, including NHB penquins in addition to QCD penguins and electrweak penguins. The motivation of this paper is based upon the fact found by us in [12], where it is found that in SUSY models, when tanβ is large, NHB are not heavy, and the mass splittings of stops and charginos are large, contributions from exchanging NHB can greatly modified the branching ratio and backward-forward asymmetry of $B \to X_s l^+ l^-, l = \mu, \tau$. It is a natural further step to explore their contributions to nonleptonic B decays. Another motivation for our work is the striking large branching

ratio for the exclusive decay of $B \to K\eta'$ [13], which has caused much theoretical efforts to assess the theoretical prediction of SM and the validity of factorization approach [14, 3, 4]. Our results show that contributions of NHB can be larger than the sum of the electroweak contributions from exchanging γ and Z bosons, and for some decay modes, they can modify both the decay widths and CP asymmetries significantly.

We shall concentrate on the decay modes in which the quark level process is $b \to q\bar{s}s, q = d,\ s$ in order to see the contributions from NHB. The transition $b \to q\bar{s}s$ is purely induced by FCNC penguin diagrams, which can be divided into four classes, 1)those of exchanging gluon, 2) of exchanging photon, 3) of exchanging Z, and 4) of exchanging Higgs neutral bosons. Normally, contributions from the first class are expected to be large, while those of the second and the third are small because of α_{em}/α_s and α_2/α_s respectively. In the SM, since suppressed by the ratio of masses of quarks and masses of NHB, those of the fourth are also negligible , so only the contributions of the first are considered. But many papers [15] (for a review, see [16]) have shown that contributions from the second and the third classes can play an important role in both the width decays and direct CP violation in some decay modes, and they can even play the major part in some processes, $B_s \to K^+K^-$, for instance. Furthermore, in SUSY context, when tanβ is large and the masses of NHB are not too heavy, contributions from exchanging neutral Higgs bosons can be considerably large and therefore can greatly modify theoretical predictions for some processes. So in the letter, we consider the case of large tanβ and take into account all contributions of the four classes.

The b transition $b \to q, q = d,\ s$; can be described by the $\Delta B = 1$ effective Hamiltonian

$$H_{\Delta B=1} = \frac{4G_F}{\sqrt{2}}[\sum_{q=u,c} V_{q'b}V^*_{q'q}(c_1O_1 + c_2O_2) - V_{tb}V^*_{tq}[\sum_{i=3}^{14} c_iO_i + \sum_{17}^{20} c_iO_i] + H.C.\ , \qquad (1)$$

where O_i are defined as

$$\begin{aligned}
O_1 &= \bar{q}_\alpha\gamma_\mu Lq'_\beta\bar{q}'_\beta\gamma^\mu Lb_\alpha \quad , \quad & O_2 &= \bar{q}\gamma_\mu Lq'\bar{q}'\gamma^\mu Lb\ , \\
O_{3(5)} &= \bar{q}\gamma_\mu Lb\sum_{q'}\bar{q}'\gamma_\mu L(R)q' \quad , \quad & O_{4(6)} &= \bar{q}_\alpha\gamma_\mu Lb_\beta\sum_{q'}\bar{q}'_\beta\gamma_\mu L(R)q'_\alpha\ , \\
O_{7(9)} &= \frac{3}{2}\bar{q}\gamma_\mu Lb\sum_{q'}e_{q'}\bar{q}'\gamma^\mu R(L)q' \quad , \quad & O_{8(10)} &= \frac{3}{2}\bar{q}_\alpha\gamma_\mu Lb_\beta\sum_{q'}e_{q'}\bar{q}'_\beta\gamma_\mu R(L)q'_\alpha\ , \\
O_{11(13)} &= \bar{q}Rb\sum_{q'}\bar{q}'L(R)q' \quad , \quad & O_{12(14)} &= \bar{q}_\alpha Rb_\beta\sum_{q'}\bar{q}'_\beta L(R)q'_\alpha\ , \\
O_{17} &= \bar{q}\sigma^{\mu\nu}Rb\sum_{q'}\bar{q}'\sigma_{\mu\nu}Rq' \quad , \quad & O_{18} &= \bar{q}_\alpha\sigma^{\mu\nu}Rb_\beta\sum_{q'}\bar{q}'_\beta\sigma_{\mu\nu}Rq'_\alpha\ , \\
O_{19} &= \frac{g_s}{8\pi^2}m_b\bar{q}\sigma^{\mu\nu}\lambda^a RbG^a_{\mu\nu} \quad . \quad & O_{20} &= \frac{e}{8\pi^2}m_b\bar{q}\sigma^{\mu\nu}RbF_{\mu\nu}
\end{aligned} \qquad (2)$$

The first two operators are induced by tree diagrams, $O_{i=3,\cdots,6}$ by gluonic penguins, $O_{i=7,\cdots,10}$ by γ and Z penguins, $O_{i=11,\cdots,14}$ and $O_{i=17,18}$ by neutral Higgs penguins, and $O_{19,20}$ by chromomagnetic and magnetic penguins[3].

[3]In order to make the minimal revisions we use the labels of operators as those in the original version of the paper. Comparing with those in refs. [26, 27, 28, 29], $O_{17}, O_{18}, O_{19}, O_{20}$ correspond to $O_{15}, O_{16}, O_{8g}, O_{7\gamma}$ respectively.

The boundary conditions of the first ten Wilson coefficients and those of $O_{19,20}$ for the solutions of the renormalization group equations are given in [6], here we provide those of $O_{11,\cdots,14,17,18}$ in the large $\tan\beta$ case.

$$
\begin{aligned}
C_{O_{11}}(m_w) &= \delta_{sq'}\frac{\alpha_2}{16\pi}\sqrt{2}tan^3\beta m_b m_s(\frac{\cos^2\alpha}{m_{H^0}^2}+\frac{\sin^2\alpha}{m_{h^0}^2}+\frac{1}{m_{A^0}^2})\left\{\sum_{i=1}^{2}\frac{m_{\chi_i}}{m_w}\right[\\
&\quad U_{i2}V_{i1}f(x_{\chi_i w},x_{\tilde{q}w})-\sum_{k=1}^{2}U_{i2}T_{k1}\,(V_{i,1}T_{k1}^* \\
&\quad -\frac{m_t}{\sqrt{2}\sin\beta}V_{i2}T_{k2}^*\Big)\,f(x_{\chi_i w},x_{\tilde{t}_k w})\Big]\Big\}. \\
C_{O_{12}}(m_w) &= 0; \\
C_{O_{13}}(m_w) &= \delta_{sq'}\frac{\alpha_2}{16\pi}\sqrt{2}tan^3\beta m_b m_s(\frac{\cos^2\alpha}{m_{H^0}^2}+\frac{\sin^2\alpha}{m_{h^0}^2}-\frac{1}{m_{A^0}^2})\left\{\sum_{i=1}^{2}\frac{m_{\chi_i}}{m_w}\right[\\
&\quad U_{i2}V_{i1}f(x_{\chi_i w},x_{\tilde{q}w})-\sum_{k=1}^{2}U_{i2}T_{k1}\,(V_{i,1}T_{k1}^* \\
&\quad -\frac{m_t}{\sqrt{2}\sin\beta}V_{i2}T_{k2}^*\Big)\,f(x_{\chi_i w},x_{\tilde{t}_k w})\Big]\Big\}. \\
C_{O_i}(m_w) &= 0,\ i=14,17,18,
\end{aligned}
\tag{3}
$$

where

$$
f(x,y)=\frac{-1}{x-y}(x\log x-y\log y),\quad x_{ij}=\frac{m_i^2}{m_j^2}
\tag{4}
$$

The conventions can be found in [12]. To reach these, we have neglected some terms contributing minor. It is interesting to mention that the terms proportional to $\tan^3\beta$ only come from the one Feynman diagram where firstly virtual NHB are emitted from b-quark line and decay into strange quark pair, then followed with flavor change caused by loops exchanging charginos and stops. While for virtual NHB decaying to down quark pair, because of the smallness of m_d/m_s the contributions are negligibly small. To up-type quark pair, SUSY contributions at most be proportional to $\tan^2\beta$ and consequently also not important compared to those for decaying to strange quark pair. This is the reason why we only pay attention to the transition $b\to q\bar{s}s$.

$C_{O_{11,13}}$ are proportional to $x_{\chi_i w}\log x_{\chi_i w}$. Therefore, contributions of SUSY depend on not only masses difference between stops, but also those between charginos. In order to escape the GIM cancellation and to increase the contributions of SUSY, both large mass splittings of stops and charginos and large chargino masses are needed. In contrast with the process $b\to s\gamma$ and the mixing of B^0 system, the smaller the mass of the lighter charginos is, the larger the contributions of SUSY.

For the first ten Wilson coefficients, we use next-leading-order QCD corrected renormalization group equations [6][4], while for the last six ones (i.e., $O_{11,\cdots,14,17,18}$), we use leading-order ones since we carried out calculations of the relevant anomalous dimension matrices only at

[4]We do not include the mixing of $O_{11,12}$ onto $O_{3,\cdots,10}$ which has been given in ref. [29].We also do not include the mixing of $O_{13,14,17,18}$ onto $O_{19,20}$ which is given first in ref. [30] and verified in ref. [29].

one-loop[5]. The anomalous dimension matrices of the last six operators (i.e., $O_{11,\cdots,14,17,18}$) can be divided into two distangled groups

$$\gamma^{(L)} = \begin{array}{c|cc} & O_{11} & O_{12} \\ \hline O_{11} & -16 & 0 \\ O_{12} & -6 & 2 \end{array} \tag{5}$$

and

$$\gamma^{(R)} = \begin{array}{c|cccc} & O_{13} & O_{14} & O_{17} & O_{18} \\ \hline O_{13} & -16 & 0 & 1/3 & -1 \\ O_{14} & -6 & 2 & -1/2 & -7/6 \\ O_{17} & 16 & -48 & 16/3 & 0 \\ O_{18} & -24 & -56 & 6 & -38/3 \end{array} \tag{6}$$

Our calculation shows that $C_{O_{12,\cdot,14,17,18}}(m_b)$ are much smaller than $C_{O_{11}}(m_b)$, so it is appropriate to neglect them in our numerical analysis.

As given in [6], the effective coefficients a_i are dependent of N_c^{eff}. In this note, we do not examine the dependence and limit ourself to a typical value of N_c^{eff}, 3, in order to see the contributions of NHBs. For $N_c^{eff} = 3$ and transition $b \to s$, the first ten effective coefficients given from Wilson coefficients in SM are [3]

$$\begin{array}{ll} a_1^{eff} = 1.05(1.05), & a_2^{eff} = 0.053(0.053), \\ a_3^{eff} = 48(48), & a_4^{eff} = -439 - 77i(-431 - 77i), \\ a_5^{eff} = -45(-45), & a_6^{eff} = -575 - 77i(-568 - 77i), \\ a_7^{eff} = 0.5 - 1.3i(0.5 - 1.3i), & a_8^{eff} = 4.6 - 0.4i(4.6 - 0.4i), \\ a_9^{eff} = -94 - 1.3i(-94 - 1.3i), & a_{10}^{eff} = -14 - 0.4i(-14 - 0.4i) \end{array} \tag{7}$$

while for transition $b \to d$, they are

$$\begin{array}{ll} a_1^{eff} = 1.05(1.05), & a_2^{eff} = 0.053(0.053), \\ a_3^{eff} = 48(48), & a_4^{eff} = -412 - 36i(-461 - 124i), \\ a_5^{eff} = -45(-45), & a_6^{eff} = -548 - 36i(-597 - 124i), \\ a_7^{eff} = 0.7 - 1i(0.3 - 1.8i), & a_8^{eff} = 4.7 - 0.3i(4.5 - 0.6i), \\ a_9^{eff} = -94 - 1i(-95 - 1.8i), & a_{10}^{eff} = -14 - 0.3i(-14 - 0.6i) \end{array} \tag{8}$$

These complex figures incorporate contributions of both the strong phases and weak phases. For the last 8 coefficients, 10^{-4} should be multiplied, and the numbers in the brackets refer to the conjugate processes $\bar{b} \to \bar{s}s\bar{s}$ and $\bar{b} \to \bar{d}s\bar{s}$.

The input parameters used in our numerical calculations are as follows.

Wolfenstein parameters [17] for the CKM matrix are given as $A = 0.81$, $\lambda = 0.2205 \pm 0.0018$ [18], $\rho = 0.12$ and $\eta = 0.34$ [19]. The scale for the running masses of quarks is set to 2.5 Gev, and the corresponding masses are given in Table I in [3]. We use the BSW approach to evaluate the form factors [20], and the form factors at zero momentum can be found in the Table II in [3]. For the reason that the form factors are not sensitive to the pole masses

[5]We would like to thank Dr. C.D. Lü for his collaboration on the calculations. A complete analysis of mixing of the operators will be published else where.

since only small extrapolations from $Q^2 = 0$ are involved in the decays $B \to h_1 h_2$, we set the pole mass to 5.4, which is given in Table III in [3], in calculating the form factors . Decay constants are given in the Table V in [3]. The mixing angles of the flavor SU(3)-octet and singlet components are set to $\theta_8 = -21.2^\circ$ and $\theta_0 = -9.2^\circ$ [21]. For more detailed information on conventions, please refer [3].

We calculate $C_{O_{11}}$ and find that it can reach 0.05. So neutral Higgs bosons can dominate the contributions of the electro-weak penguins for some processes. In calculations the relevant ten supersymmetric parameters are taken as

$$\begin{array}{llll} M_2 = 300\ GeV, & \mu = 300\ GeV, & m_{A^0} = 80\ GeV, & m_{\tilde{u}} = 500\ GeV, \\ m_R = 300\ GeV, & m_L = 300\ GeV, & A_t = 300 GeV, & tan\beta = 50 \\ & \psi_\mu = \pi & \psi_{A_t} = 0.3 & , \end{array} \tag{9}$$

which satisfy the constraint on the SUSY phases arising from EDMs through the Barr-Zee mechanism [22, 23] as well as the constraint from $b \to s\gamma$ [12].

we get $C_{O_{11}}(m_b) = -0.051 - 0.027i$, and consequently $a_{11}^{eff} = -0.051 - 0.027i$ and $a_{12}^{eff} = -0.017 - 0.009i$. We note that the parameters (9) satisfy the constraint from $b \to s\gamma$. For the first ten Wilson coefficients, we find that the contributions of SUSY with the above parameters can be safely neglected, because they only modify the Wilson coefificients in few percent in most of part of the parameter space. So we shall still use eqs. (7) and (8) for the sake of simplicity. The complex phase of $C_{O_{11}}$ is originate from the SUSY CP phases, if $\psi_{A_t} = 0$, the corresponding $a_{11}^{eff} = -0.055$, no imaginary part appears.

To evaluate the relevant transition matrix elements, we will use the naive factorization assumption. In the spectator model, several formula on decay amplitudes are listed below

$$\begin{aligned} M(\bar{B}^0 \to \bar{K}^0 \eta^{(\prime)}) &= -\frac{G_F}{\sqrt{2}} M_{sss}^{B\to\eta^{(\prime)}}(K) V_{tb} V_{ts}^* \Big[a_4 - \frac{1}{2} a_{10} + \\ &\quad (a_6 - \frac{1}{2} a_8) R_5\Big] - \frac{G_F}{\sqrt{2}} M_{sss}^{B\to K} \{V_{ub} V_{us}^* a_1 + \\ &\quad V_{cb} V_{cs}^* a_1 \frac{f_{\eta^{(\prime)}}^c}{f_{\eta^{(\prime)}}^u} - V_{tb} V_{ts}^* \Big[2a_3 - 2a_5 + \frac{1}{2}(a_9 - a_7) + (a_3 - a_5 \\ &\quad + a_9 - a_7) \frac{f_{\eta^{(\prime)}}^c}{f_{\eta^{(\prime)}}^u} + \Big(a_3 + a_4 - a_5 + \frac{1}{2}(a_7 - a_9 - a_{10} - a_{12}) + \\ &\quad (a_6 - \frac{1}{2}(a_8 - a_{11})) R_6^{(\prime)}\Big) (\frac{f_{\eta^{(\prime)}}^s}{f_{\eta^{(\prime)}}^u} - 1)\Big]\Big\}. \\ M(\bar{B}^0 \to \bar{K}^0 \phi) &= -\sqrt{2} G_F M_{sss}^{B\to K}(\phi) V_{tb} V_{ts}^* \big[a_3 + a_4 + a_5 \\ &\quad - \frac{1}{2}(a_7 + a_9 + a_{10} + a_{12})\big] \end{aligned} \tag{10}$$

Where

$$\begin{aligned} M_{q_1 q_2 q_3}^{B\to P}(P'(V)) &= \langle P'(V)|(\bar{q}_1 q_2)_{V-A}|0\rangle \langle P|(\bar{q}_3 b)_{V-A}|\bar{B}^0\rangle, \\ R_5 &= \frac{2m_K^2}{(m_b - m_u)(m_u + m_s)}, \\ R_6^{(\prime)} &= \frac{2m_{\eta^{(\prime)}}}{(m_b - m_s)(m_s + m_s)}. \end{aligned} \tag{11}$$

and the superscript "eff" of a_i^{eff} has been omitted for the sake of simplicity.

Compared with formula given in [3], we have taken into account the contributions of NHB's by adding a_{11}, a_{12} and a_{13} in these amplitudes of decays. We have not included W-exchange, W-annihilation and spacelike penguin contributions because the W-exchange and W-annihilation contributions are negligible due to helicity suppression (as well as color suppression for W-exchange) and we do not have a reliable method to estimate the spacelike penguins[4]. We have used the equations of motion and Fierz rearrangement to calculate the contributions of scalar operators. For other formula, it is easy to derive when compared with the tables given in the appendices in [3, 4]. Our numerical analysis about the effects of NHB are based on these formula.

For charged $B^{\pm}$, A_{CP} is defined as

$$A_{CP} = \frac{\Gamma(B^- \to f^-) - \Gamma(B^+ \to f^+)}{\Gamma(B^- \to f^-) + \Gamma(B^+ \to f^+)} \tag{12}$$

and it reflects the magnitude of the direct CP violation. For the cases in which both B^0 and $\bar{B}^0$ decay to the same final states, A_{CP} is defined as

$$A_{CP} = \frac{1 - |\epsilon|^2 - 2Im(\epsilon x_B)}{(1 + |\epsilon|^2)(1 + x_B^2)} \tag{13}$$

where $\epsilon = \frac{q}{p}\frac{\bar{A}}{A}$, a quantity reflecting effects from both the mixing and the direct CP violations. We omit the analysis on the CP violation for decays, $B^0 \to VV'$, and $\bar{B}^0 \to \bar{K}^{*0}K^0, K^{*0}\bar{K}^0$, in order to simplify the calculations. x_B is set to 0.723 as given in [24]. We do not introduce new CP origin in the mixing of the neutral B meson.

Our result is shown in Table 1. A number of observations are in order.

1. Compared to the SM, the decay widths for most modes corresponding to the transition $b \to q\bar{s}s$ increase. For example, for $B^- \to K^-\eta'$, contributions of NHB increase the branching ratio up to 25% and for $B^- \to K^-\eta$, they increase up to 80%. For $B \to K(K^*)\phi$ ($\bar{B}^0 \to \pi^0\eta'$) because of the cancellations between $a_{12}(a_{11})$ and $a_4(a_6)$, the decay widths decrease. We see that the contributions of NHB do modify the branching ratios for all charmless decay modes whose amplitudes are governed by penguins for transition $b \to q\bar{s}s$, though for most nonleptonic decays gluon penguins dominate. Compared with the semileptonic decays $b \to sl^+l^-$, where branching ratios and forward-backward asymmetry can be greatly modified [12] by the effects of NHB, for most charmless nonleptonic decay modes their effects are soften by the large values of the coefficients a_1, a_4, a_6.

2. For those decay modes where electroweak penguins play an important role, such as $B \to K^*\eta$, $\bar{B}^0 \to \bar{K}^{*0}(K^0)\eta$, $\bar{B}^0 \to \pi^0\eta$,$B^0 \to \eta(\eta')\eta(\eta')$, and $\bar{B}^0 \to \omega^0\eta'$, the contributions of NHB can enhance the branching ratios by several handreds percent.

3. For the modes where there are subtle cancellations between contributions in SM, such as $B \to K^{*0}K$, $\bar{B}^0 \to K^{*0}\eta'$, and $\bar{B}^0 \to \rho^0\eta'$, the contributions of NHB can enhance the branching ratios by a factor of a few tens and greatly modify the corresponding CP asymmetry.

In the effective Hamiltonian method to study hadronic decays the largest theoretical uncertainty arises from the calculations of hadronic matrix elements of local operators. For energetic two-body decays of B mesons factorization should be a good approximation because of color transparency [25] and non-factorization effects may be parameterized by N_c^{eff}.

In the factorization framework, in addition to the uncertainty due to non-factorization effects, the other uncertainties come from the input parameters, such as form factors, decay constants, quark masses, CKM matrix elements, the scale at which Wilcon coefficients is calculated, etc. The different choices of input parameters can lead to an uncertainty of from a few percent to several handreds percent (dependent on modes), as can be seen by comparing the results for the same N_c^{eff} between ref.[3] and ref.[4]. The uncertainty due to different N_c^{eff} can also reach several handreds percent. Therefore, it is difficult to disentagle the non-standard physics for most of modes in Table 1. However, by comparing our result and the results given in [3] and [4], it can be drawn that
A. compared to SM, a significantly large enhancement of the branching retio for $B \to K^*\eta'$ in MSSM is definitely existed; and
B. the branching ratio for $B \to \rho^0\eta'$ ($\bar{B}^0 \to K^{*0}\bar{K}^0$) is enhanced by a factor of 40 (25) compared to SM for the same N_c^{eff} and larger than all values corresponding to different N_c^{eff} in SM given in [3, 4]. Thus, it could be possible to disentagle MSSM from SM by measuring the branching ratio if tanβ is large and NHBs are not heavy.

In summary, we have analyzed the contributions of NHB to charmless nonleptonic two-body B decays. Our analysis shows that at some regions of the parameter space in MSSM, NHBs can dominate the contributions of electroweak penguins and significantly modify both decay widths and CP asymmetry of the decay modes which are due to the transition $b \to q\bar{s}s$. Our analysis has not include CP phase of SUSY origin in mixing. After taking into the complex contributions of SUSY to the mixing, CP asymmetry of most of neutral B meson decay modes could be modified.

¿From the above analysis, we know that the effects of NHB, for most charmless nonleptonic decay modes, are soften by the large values of the coefficients a_1, a_4, a_6. But their effects can show up when there are subtle cancellations. Therefore, although there are huge theoretical uncertainties, the enhancements for some modes, such as $B \to \rho^0\eta'$ and $\bar{B}^0 \to K^{*0}\bar{K}^0$, exceed the theoretical uncertainties significantly so that it would be possible to search new physics effects by measuring these decay modes.

Acknowledgment

C.S. Huang would like to acknowledge KIAS where part of the letter was written for the warm hospitality. Q.S. Yan would like to thank the helpful discussions with Dr. O. Vives, Dr. C.D. Lü and Dr. K.C. Yang. This work was supported in part by the National Natural Science Foundation of China and partly supported by CCAST.

References

[1] B.H. Behrens *et al.*, *Phys. Rev. Lett.* **80** (1998) 3710; T. Bergfeld *et al.*, *Phys. Rev. Lett.* **81** (1998) 272; T.E. Browder *et al.*, *Phys. Rev. Lett.* **81** (1998) 1786.

[2] H.Y. Cheng, hep-ph/9902378.

[3] A. Ali, G. Kramer, and C.D. Lü, *Phys. Rev.* **D 58** (1998) 094009; *Phys. Rev.* **D 59** (1998) 014005.

[4] Y.H. Chen, H.Y. Cheng, B.Tseng, and K.C. Yang, hep-ph/9903453; H.Y. Cheng and B. Tseng, *Phys. Rev.* **D 58** (1998) 094005.

[5] D. Atwood, A. Soni, *Phys. Lett.* **B 405** (1997) 150; W.S. Hou and B. Tseng, *Phys. Rev. Lett.* **80** (1998) 434; A.L. Kagan, hep-ph/9806266; H. Fritzsch, *Phys. Lett.* **B 415** (1997) 83; X.G. He, W.S. Hou, and C.S. Huang, *Phys. Lett.* **B 429** (1998) 99; X.G. He, *Phys. Lett.* **B 454** (1999) 123, and references therein.

[6] G. Buchalla, A.J. Buras, and M.E. Lauthenbache, *Rev. Mod. Phys.* **68** (1996) 1125; A. Buras, *et al.*, *Nucl. Phys.* **B 400** (1993) 37; *Nucl. Phys.* **B 400** (1993) 75; S. Bertonlini, *et al.*, *Nucl. Phys.* **B 353** (1991) 591.

[7] W.N. Cottingham, H. Mehrban, and I.B. Whittingham, hep-ph/9905300.

[8] H. Könic, *Zeit. für Physik* **C 73** (1996) 161.

[9] D. Choudhury, B. Dutta, and A. Kundu, hep-ph/9812209; D. Guetta, hep-ph/9805274.

[10] F. Gabbiani, E. Gabrielli, A. Masiero, and L.Silvestrini, *Nucl. Phys.* **B 477** (1996) 321; A. Masiero, and L. Silvestrini, hep-ph/9711401; hep-ph/9709242; hep-ph/9709244.

[11] R. Barbieri, and A. Strumia, *Nucl. Phys.* **B 508** (1997) 3; S.A. Abel, W.N.Cottingham, and I.B. Whittingham, *Phys. Rev.* **D 58** (1998) 073006; W.N. Cottingham, H.Mehrban, and I.B. Whittingham, hep-ph/9905300; D.S. Du and M.Z. Yang, J. Phys. C: Nucl. Part. Phys. 24(1998)2009. F. Krauss, and G. Soff, hep-ph/9807238; S. Bertolini, F. Borzynatu, A. Masiero and G. Ridolfi, *Nucl. Phys.* **B 353** (1991) 591; G.C. Branco, G.C. Cho, Y. Kizukuri, and N. Oshimo, *Nucl. Phys.* **B 449** (1995) 483; *Phys. Lett.* **B 337** (1994) 316; T. Goto, T. Nihei, and Y. Okada, *Phys. Rev.* **D 53** (1996) 5233; G. Barenboim, and M. Raidal, hep-ph/9903270. T. Goto, Y. Okada and Y. Shimizu, *Phys. Rev.* **D 58** (1998) 094006, and references therein; B. Lunghi, A. Masiero, I. Scimemi, and L. Silvestrini, hep-ph/9906286.

[12] C.S. Huang and Q.S. Yan, *Phys. Lett.* **B 442** (1998) 209; C.S. Huang, W. Liao and Q.S. Yan, *Phys. Rev.* **D 59** (1998) 011701.

[13] CLEO collaboration, B.H. Behrens et al, *Phys. Rev. Lett.* **80** (1998) 3710; J.G. Smith, hep-ex/9803028.

[14] X.G. He and G. L. Lin, hep-ph/9812248; D.S. Du, Y.D. Yang, and G.H Zhu, *Phys. Rev.* **D 59** (1999) 014007.

[15] N.G. Despande and X.G. He, *Phys. Rev. Lett.* **74** (1995) 26; N.G. Despande, X.G. He and J. Trampetic, *Phys. Lett.* **B 345** (1995) 547; R. Fleischer, *Phys. Lett.* **B 365** (1996) 399; D.S. Du and M.Z. Yang, *Phys. Lett.* **B 358** (1995) 123.

[16] M. Gronau, hep-ph/9904329;

[17] L. Wolfenstein, *Phys. Rev. Lett.* **51** (1983) 1945.

[18] A. Ali, DESY Report 97-256, hep-ph/9801270; to be published in Proc. of the First APCTP Workshop, Pacific Particle Physics Phenomenology, Seoul, South Korea. For further details, see A. Ali and D. London, Nucl. Phys. B (Proc. Suppl.) **54A**, 297 (1997).

[19] R.M. Barnett et al. (Particle Data Group), Phys. Rev. **D54**, 1 (1996).

[20] M. Bauer and B. Stech, Phys. Lett. **B152**, 380 (1985);
M. Bauer, B. Stech and M. Wirbel, Z. Phys. **C34**, 103 (1987).

[21] T. Feldmann, P. Kroll, and B. Stech, *Phys. Rev.* **D 58** (1998) 114006; hep-ph/9812269. M.A. Diaz, *Phys. Lett.* **B 322** (1994) 591; T. Goto and Y. Okada, Prog. of Theor. Phys. 94 (1995) 407; R. Garisto and J.N. Ng, *Phys. Lett.* **B 315** (1993) 372; J.L. Hewett, SLAC-PUB-6521 May 1994 T/E; F.M. Borzumati, M. Olechowski and S. Pokorski, CERN-TH 7515/94, TUM-73-83/94.

[22] S.M.Barr and A.Zee, Phys.Rev.Lett.**65**, 21(1990)

[23] D. Chang, W.-Y. Keung, A. Pilaftsis, Phys. Rev. Lett. **82** (1999) 900.

[24] Particle Data Group, E. P. J. C3(98)1.

[25] J. D. Bjorken, Nucl. Phys. **B** (Proc. Suppl.) **11** (1989) 325.

[26] J.-F. Cheng, C.-S. Huang, Phys. Lett. **B554** (2003) 155.

[27] J.-F. Cheng, C.-S. Huang, and X.-H. Wu, hep-ph/0306086.

[28] C.-S. Huang and S.-H. Zhu, hep-ph/0307354, to appear in Phys. Rev. D.

[29] G. Hiller and F. Krüger, hep-ph/0310219.

[30] F. Borzumati, C. Greub, T. Hurth and D. Wyler, Phys. Rev. **D62**, 075005 (2000).

Table 1: Branching ratios in unit of 10^{-6},and CP asymmetry in unit of %.

$N_c = 3$	Branching Ratio		CP Asymmetry	
Decay Modes	SM	SUSY	SM	SUSY
$B^- \to K^- \eta$	1.8	3.27	-6.1	-0.16
$B^- \to K^- \eta'$	25.2	31.9	4.5	5.58
$B^- \to K^- \phi$	13.8	8.11	2.0	-4.8
$B^- \to K^{*-} \eta$	1.50	6.64	5.4	3.3
$B^- \to K^{*-} \eta'$	0.89	6.71	23.7	-5.4
$B^- \to K^{*-} \phi$	6.95	4.43	2.0	-3.7
$B^- \to K^- K^0$	0.67	0.82	-12.3	-9.8
$B^- \to \pi^- \eta$	2.10	2.14	-14.0	-13.3
$B^- \to \pi^- \eta'$	2.16	2.56	-13.0	-11.7
$B^- \to K^0 K^{*-}$	0.46	0.69	-14.5	-9.2
$B^- \to K^{*0} K^-$	0.0004	0.01	-46.7	6.2
$B^- \to \rho^- \eta$	6.16	6.89	-3.4	-3.1
$B^- \to \rho^- \eta'$	7.41	6.99	-11.0	-11.3
$B^- \to K^{*-} K^{*0}$	5.73	3.77	-14.5	-22.1
$\bar{B}^0 \to \bar{K}^0 \eta$	1.6	3.1	27.8	37.5
$\bar{B}^0 \to \bar{K}^0 \eta'$	24.1	30.5	29.4	33.5
$\bar{B}^0 \to \phi \bar{K}^0$	13.3	7.8	29.3	15.1
$\bar{B}^0 \to \bar{K}^{*0} \eta$	1.43	6.14	30.0	41.4
$\bar{B}^0 \to \bar{K}^{*0} \eta'$	0.43	6.4	24.6	37.2
$\bar{B}^0 \to \bar{K}^{*0} \phi$	6.69	4.27	—	—
$\bar{B}^0 \to \bar{K}^0 K^0$	0.65	0.79	33.2	35.6
$\bar{B}^0 \to \pi^0 \eta$	0.32	0.71	31.3	38.8
$\bar{B}^0 \to \pi^0 \eta'$	0.12	0.019	26.8	-70.3
$\bar{B}^0 \to \eta \eta'$	0.039	0.11	24.0	39.0
$\bar{B}^0 \to \eta \eta$	0.13	0.35	24.0	39.0
$\bar{B}^0 \to \eta' \eta'$	0.024	0.08	-4.3	-27.2
$\bar{B}^0 \to K^0 \bar{K}^{*0}$	0.44	0.66	—	—
$\bar{B}^0 \to K^{*0} \bar{K}^0$	0.0004	0.01	—	—
$\bar{B}^0 \to \rho \eta$	0.015	0.0076	34.2	-34.7
$\bar{B}^0 \to \rho \eta'$	0.0028	0.11	63.1	27.3
$\bar{B}^0 \to \omega \eta$	0.046	0.048	-10.5	-56.3
$\bar{B}^0 \to \omega \eta'$	0.019	0.11	-57.9	36.6
$\bar{B}^0 \to \bar{K}^{*0} K^{*0}$	5.52	3.63	—	—